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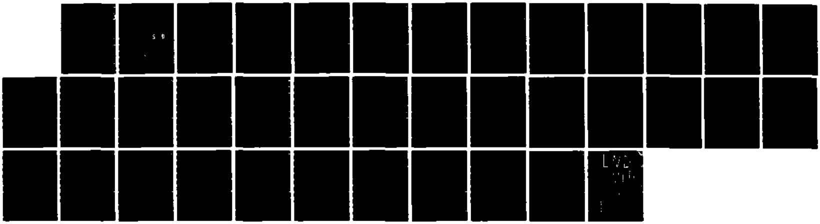

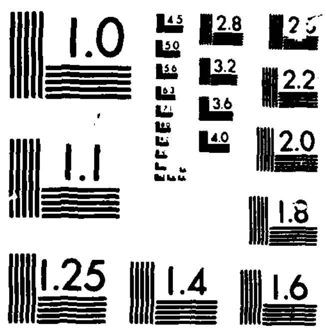

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The Effect of Finite "Blob" Size on the Current Convective Instability in the Auroral Ionosphere

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April 11, 1986

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It has been suggested that the current convective instability may be responsible for the structuring, i.e., generation of density irregularities, of density enhancements (known as "blobs") in the auroral ionosphere. However, previous theories have neglected the finite extent of the "blob" along the geomagnetic field. In this paper we develop a nonlocal theory of the current convective instability which considers the finite extent of an ionospheric "blob" parallel to the geomagnetic field. We find that the growth rate of the instability can be substantially reduced in the finite-sized "blob" case from the value obtained in the local approximation for an infinitely long blob. For auroral ionosphere parameters, the reduction in the growth rate for medium scale irregularities (1-10 km) can be one to two orders of magnitude for the typical observed values of "blob" sizes (~ a few hundred km). Thus, it appears that the current convective instability is not a viable mechanism to generate scintillation causing irregularities, i.e., 1-10 km irregularities.

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CONTENTS

THE EFFECT OF FINITE "BLOB" SIZE ON THE CURRENT CONVECTIVE INSTABILITY IN THE AURORAL IONOSPHERE

I. INTRODUCTION

Ionospheric irregularities (or density fluctuations) cause scintillation of radio signals which has been observed during diffuse auroral situations [Fremouw et al., 1977; Rino et al., 1978; Rino and Owen, 1980]. The irregularities occur in regions of soft particle precipitation (which are confined in latitude) and horizontal plasma density gradients. At present, the emerging picture is that large-scale density enhancements (known as "blobs") are produced in the high latitude F region and are convected around the polar ionosphere; the sides of these "blobs" appear to be the regions of structure [Tsunoda and Vickrey, 1985]. There has been a considerable theoretical effort to understand these irregularities in terms of gradient driven fluid instabilities [Ossakow and Chaturvedi, 1979; Huba and Ossakow, 1980; Vickrey et al., 1980; Keskinen et al., 1980; Chaturvedi and Ossakow, 1981, 1983; Keskinen and Ossakow, 1983; Satyanarayana and Ossakow, 1983; Gary, 1984; Huba, 1984; Satyanarayana et al., 1985]. In particular, the current convective instability [Lehnert, 1958; Kadomtsev and Nedospasov, 1960] has been suggested as a generation mechanism of plasma irregularities in the high latitude ionosphere [Ossakow and Chaturvedi, 1979; Fejer and Kelley, 1980; Hanuise et al., 1981; Vickrey and Kelley, 1983]. This instability can be excited in a weakly collisional, magnetized plasma which contains a field-aligned current and a transverse density gradient. This configuration exists, at times, in the auroral ionosphere along the sides of density enhancements ("blobs").

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The theoretical study of the current convective instability, as applied to the auroral F region, has evolved over the years. Initial studies used the local approximation which is valid for modes such that kL >> 1 where k is the wavenumber and L is the scale length of the density gradient [Ossakow and Chaturvedi, 1979; Vickrey et al., 1980; Chaturvedi and Ossakow, 1981; Gary, 1984]. Subsequent studies have considered a host of nonlocal effects which can be important and are neglected in a local analysis, e.g., magnetic shear [Huba and Ossakow, 1980]; velocity shear [Satyanarayana and Ossakow, 1983]; long wavelength modes (i.e., kL < 1) [Huba, 1984]; finite current channel width [Satyanarayana et al., 1985]. Overall, these effects have a tendency to reduce the growth rate of the instability, with the exception of velocity shear which can stabilize short wavelength modes (those with kL >> 1). Despite the modest reductions in growth rate the current convective instability has remained a viable mechanism to generate density irregularities with scale lengths 1-10 km.

A potentially important nonlocal effect that has not been considered to date is the finite extent of the "blob" along the ambient geomagnetic field. For example, recent studies of barium cloud dynamics have indicated that the parallel extent of the cloud along the field can affect its stability properties [Goldman et al., 1976; Sperling et al., 1984; Sperling and Glassman, 1985; Drake et al., 1985]. Physically, the reason that this effect could be significant is the following. The current convective instability requires density and potential fluctuations both parallel and perpendicular to $\underline{B}_0$, i.e., assuming a Fourier expansion of the modes we have $\underline{k} = k_\perp \hat{e}_\perp + k_\parallel \hat{e}_\parallel$ where $\underline{k}$ is the wavenumber. Based on local theory, the instability attains maximum growth when $k_\perp/k_\parallel \sim (\sigma_\parallel/\sigma_\perp)^{1/2}$ where $\sigma_\parallel$ and $\sigma_\perp$ are the plasma conductivities parallel and perpendicular to the

field, respectively. For typical F region parameters we note that $\sigma_{\parallel}/\sigma_{\perp} \sim 10^8$ so that $\lambda_{\parallel} \sim 10^4\, \lambda_{\perp}$ where λ is the wavelength of the mode. Thus, for fluctuations with $\lambda_{\perp} \sim 1$-10 km, local theory requires parallel wavelengths such that $\lambda_{\parallel} \sim 10^4 - 10^5$ km for maximum growth. On the other hand, the parallel extent of a "blob" is typically only several hundred km so that local theory is probably inadequate to properly describe the current convective instability in the auroral ionosphere (at least for the fastest growing modes of interest).

In this paper we develop a nonlocal theory of the current convective instability which considers the finite extent of an ionospheric "blob" parallel to the geomagnetic field. We show that the physical picture presented in the previous paragraph is, in fact, reasonably accurate. For a mode that has a parallel structure sufficiently short to "fit" the fastest growing mode into the "blob", we recover the maximum growth rate predicted by local theory. However, modes that have a parallel structure determined by the extent of the "blob" along the field can have growth rates substantially reduced from the maximum growth rate expected from local theory. For typical auroral ionosphere parameters, the reduction in the maximum growth rate of the instability for medium scale irregularities (1-10 km) can be one to two orders of magnitude.

The organization of the paper is as follows. In the next section we present the physical model and basic equations used in the analysis. In Section III we derive the dispersion equation for the current convective instability which explicitly includes the finite extent of the "blob" along the ambient magnetic field. In Section IV we present analytical and numerical results. Finally, in Section V we summarize our findings and apply our results to the auroral ionosphere.

II. MODEL AND FUNDAMENTAL EQUATIONS

We consider a slab geometry and a plasma configuration as shown in Fig. 1. The ambient magnetic field and current are in the z-direction ($\underset{\sim}{B} = B_0 \hat{e}_z$ and $\underset{\sim}{J} = J_0 \hat{e}_z$). The ambient density profile used is given by

$$n_0(x,z) = \begin{cases} n_b(x) & 0 < z < z_0 \\ n_i & \text{otherwise} \end{cases} \tag{1}$$

where n_b denotes the "blob" density which is confined to a limited region along $\underset{\sim}{B}_0$ and is inhomogeneous transverse to $\underset{\sim}{B}_0$, and n_i denotes the homogeneous background ionosphere. We consider low frequency fluctuations such that $\partial/\partial t \ll \Omega_\alpha, \nu_{\alpha n}$ where Ω_α and $\nu_{\alpha n}$ represent the gyrofrequency and collision frequency with neutrals associated with the α species. We also assume $\nu_{en}/\Omega_e \ll \nu_{in}/\Omega_i \ll 1$ which is relevant to the F region of the ionosphere. Finally, for simplicity we assume a cold plasma (i.e., $T_e = T_i = 0$) and that the ambient current is carried by electrons (i.e., $\underset{\sim}{J}_0 = - n_e V_{0z}\hat{e}_z$).

Within the context of our assumptions, the basic equations of our analysis are continuity, momentum transfer, charge neutrality, and Ampere's law:

$$\frac{\partial n_\alpha}{\partial t} + \nabla\cdot\left(n_\alpha \underset{\sim}{v}_\alpha\right) = 0 \tag{2}$$

$$0 = - e\underset{\sim}{E} - \frac{e}{c}\underset{\sim}{v}_e \times \underset{\sim}{B} - m_e \nu_{en} \underset{\sim}{v}_e \tag{3}$$

$$0 = e\underset{\sim}{E} + \frac{e}{c}\underset{\sim}{v}_i \times \underset{\sim}{B} - m_i \nu_{in} \underset{\sim}{v}_i \tag{4}$$

$$\nabla\cdot\underset{\sim}{J} = \nabla\cdot[ne(\underset{\sim}{v}_i - \underset{\sim}{v}_e)] = 0 \tag{5}$$

$$\nabla \times \underset{\sim}{B} = \frac{4\pi}{c}\underset{\sim}{J} \tag{6}$$

where the variables have their usual meanings and we are in the neutral frame of reference (i.e., $\underset{\sim}{V}_n = 0$). We take the electric and magnetic fields to be represented by potentials as

$$\underset{\sim}{E} = -\nabla\phi - \frac{1}{c}\frac{\partial A_z}{\partial t}\hat{e}_z \tag{7}$$

and

$$\underset{\sim}{B} = B_0\hat{e}_z + \nabla A_z \times \hat{e}_z \tag{8}$$

where B_0 is the ambient field, and ϕ and A_z are the electrostatic and vector potentials, respectively. We consider only A_z since $J_\parallel \gg J_\perp$.

The electron cross-field motion is given by

$$\underset{\sim}{v}_{e\perp} = -\frac{c}{B}\nabla_\perp\phi \times \hat{e}_z \tag{9}$$

while the parallel motion is given by

$$v_{e\parallel} = \frac{e}{m_e\nu_{en}}\left[\nabla_\parallel\phi + \frac{1}{c}\frac{\partial A_z}{\partial t}\right]. \tag{10}$$

The ion cross-field motion is given by

$$\underset{\sim}{v}_{i\perp} = - \frac{c}{B} \nabla_\perp \phi \times \hat{e}_z - \frac{\nu_{in}}{\Omega_i} \frac{c}{B} \nabla_\perp \phi \tag{11}$$

and the parallel motion is given by

$$v_{i\parallel} = - \frac{m_e}{m_i} \frac{\nu_{en}}{\nu_{in}} v_{e\parallel} << v_{e\parallel}. \tag{12}$$

We now substitute (9)-(12) into (2), (5), and (6) and arrive at the following equations:

$$\frac{\partial n}{\partial t} - \frac{c}{B} \nabla_\perp \phi \times \hat{e}_z \cdot \nabla n + \frac{}{4\pi e} \frac{\partial}{\partial z} \nabla_\perp^2 A_z = 0 \tag{13}$$

$$\frac{\nu_{in}}{\Omega_i} \frac{c}{B} \nabla \cdot \left(n \nabla_\perp \phi \right) + \frac{c}{4\pi e} \frac{\partial}{\partial z} \nabla_\perp^2 A_z = 0 \tag{14}$$

$$\nabla_\perp^2 A_z = \frac{4\pi}{c\eta_e} \frac{\partial \phi}{\partial z} + \frac{1}{c} \frac{\partial A_z}{\partial t} \tag{15}$$

where $\eta_e = m_e \nu_{en} / ne^2$.

The equilibrium solution of (13)-(15) requires that

$$- \frac{c}{B} \nabla_\perp \phi_0 \times \hat{e}_z \cdot \nabla n_0 + \frac{c}{4\pi e} \frac{\partial}{\partial z} \nabla_\perp^2 A_{z0} = 0 \tag{16}$$

$$\frac{\nu_{in}}{\Omega_i} \frac{c}{B} \nabla \cdot \left(n_0 \nabla_\perp \phi_0 \right) + \frac{c}{4\pi e} \frac{\partial}{\partial z} \nabla_\perp^2 A_{z0} = 0 \tag{17}$$

$$\nabla_\perp^2 A_{z0} = \frac{4\pi}{c\eta_e} \frac{\partial \phi_0}{\partial z} \tag{18}$$

We have assumed $n_0 = n_0(x,z)$ [see (1)] and will also take ϕ_0 and A_{z0} to be functions only of x and z. [Neglect of the y dependence on ϕ_0 is equivalent to assuming $E_{0y} = 0$; we note that it is the y component

of $\underline{E}_0$ which can drive the usual $\underline{E} \times \underline{B}$ gradient drift instability in the F region and we are excluding this possibility.] From (16) and (17) we then find that

$$\frac{c}{4\pi e} \frac{\partial}{\partial z} \nabla_\perp^2 A_{z0} = 0 \tag{19}$$

and

$$\frac{\nu_{in}}{\Omega_i} \frac{c}{B} \frac{\partial}{\partial x} \left(n_0 \nabla_\perp \phi_0\right) = 0 \tag{20}$$

Substituting (18) into (19) we find that

$$\frac{\partial}{\partial z} \frac{4\pi}{c n_e} \frac{\partial \phi_0}{\partial z} = 0 \tag{21}$$

Making use of the definition of n_e $(= m_e \nu_{en}/ne^2)$ and using (20) and (21) we can rewrite (14) as

$$\frac{\nu_{in}}{\Omega_i} \frac{\partial}{\partial x} \left(n_0 \frac{\partial \phi_0}{\partial x}\right) + \frac{\partial}{\partial z} \left(\frac{\Omega_e}{\nu_{en}} n_0 \frac{\partial \phi_0}{\partial z}\right) = 0 \tag{22}$$

which is equivalent to the requirement that $\nabla \cdot \underline{J}_0 = 0$ since $J_{0x} = - n_0 (\nu_{in}/\Omega_i)(c/B)\partial\phi_0/\partial x$ and $J_{0z} = (n_0 e/m_e \nu_{en})\partial\phi_0/\partial z$.

Finally, we can write an equilibrium potential to be

$$\phi_0(x,z) = \begin{cases} - E_{0z} z - E_{0x} \int^x \dfrac{n_i \, dx'}{n_b(x')} & 0 < z < z_0 \\ - E'_{0z} z & \text{otherwise} \end{cases} \tag{23}$$

This potential generates an equilibrium electric field

$$\mathbf{E} = \begin{cases} E_{0z}\,\hat{e}_z + E_{0x} n_i/n_b(x)\,\hat{e}_x & 0 < z < z_0 \\ E'_{0z}\,\hat{e}_z & \text{otherwise} \end{cases} \tag{24}$$

which causes a parallel electron drift $\underset{\sim}{V}_{0z} = -\left(eE_{0z}/m_e\nu_{en}\right)\hat{e}_z$ and an inhomogeneous cross-field drift $\underset{\sim}{V}_{0\perp}(x) = -\left(cE_{0x}/B\right)\left(n_i/n_b(x)\right)\hat{e}_y$ in the "blob" plasma. In the background plasma there is only a parallel electron drift $\underset{\sim}{V}'_{0z} = -\left(eE'_{0z}/m_e\nu_{en}\right)\hat{e}_z$.

III. DISPERSION EQUATION

We linearize (13)-(15) in order to obtain the dispersion equation which describes the current convective instability. We use the equilibrium derived in Section II and assume perturbed quantities vary as $\tilde{p} = \tilde{p}(x)\exp(\gamma t - ik_y y)$. The linearized equations are

$$\frac{\partial\tilde{n}}{\partial t} - \frac{c}{B}\nabla_\perp\tilde{\phi}\times\hat{e}_z\cdot\nabla_\perp n_0 - \frac{c}{B}\nabla_\perp\phi_0\times\hat{e}_z\cdot\nabla\tilde{n} + \frac{c}{4\pi e}\frac{\partial}{\partial z}\nabla_\perp^2\tilde{A}_z = 0 \tag{25}$$

$$\frac{\nu_{in}}{\Omega_i}\frac{c}{B}\nabla\cdot\left(n_0\nabla_\perp\tilde{\phi}\right) + \frac{\nu_{in}}{\Omega_i}\frac{c}{B}\nabla\cdot\left(\tilde{n}\nabla_\perp\phi_0\right) + \frac{c}{4\pi e}\frac{\partial}{\partial z}\nabla_\perp^2\tilde{A}_z = 0 \tag{26}$$

and

$$\nabla_\perp^2\tilde{A}_z = \frac{4\pi e}{c}\tilde{n}V_{0z} + \frac{4\pi}{cn_e}\frac{\partial\tilde{\phi}}{\partial z} + \frac{1}{c}\frac{\partial\tilde{A}_z}{\partial t}\,. \tag{27}$$

In the following we will ignore terms proportional to $\nabla_\perp\phi_0$ which is valid for $k_y V_{0\perp} \ll |V_{0z}\,\partial/\partial z|$. Retaining these terms leads to two effects: (1)

a real frequency is produced proportional to $k_y V_{0\perp}(x_0)$ where $x = x_0$ is the location of mode localization, and (2) velocity shear stabilization of short wavelength modes (Satyanarayana and Ossakow, 1983). We are intentionally ignoring these effects in order to highlight the influence of finite "blob" size on the current convective instability.

Eliminating $\tilde{n}$ from these equations we obtain two coupled differential equations for $\tilde{A}_z$ and $\tilde{\phi}$,

$$[\gamma + k_y^2 D_r]\ \tilde{A}_z = - c\ \frac{\partial\tilde{\phi}}{\partial z} + i\ \frac{\gamma_0}{\gamma\sqrt{\alpha}}\ ck_y\tilde{\phi} \tag{28}$$

and

$$\tilde{\phi} = - \alpha\ D_r\ \frac{1}{c}\ \frac{\partial\tilde{A}_z}{\partial z} \tag{29}$$

where $D_r = (c^2/\omega_{pe}^2)\nu_e$ is the resistive diffusion coefficient, $\omega_{pe}^2 = 4\pi n_0 e^2/m_e$, $\alpha = (\Omega_e/\nu_e)(\Omega_i/\nu_{in})$, and $\gamma_0 = -(\nu_e/\Omega_e)\sqrt{\alpha}\ V_{0z}\ \partial\ell n\ n_0/\partial x$. In writing (28) we have taken $(V_{0z}/\gamma)\partial/\partial z << 1$ which can be shown a posteriori; this term does not affect the growth rate of the mode but only causes the mode to have a real frequency. Finally, combining (28) and (29), we obtain the mode equation for the current convective instability in terms of the vector potential $\tilde{A}_z$,

$$\frac{\partial}{\partial z}\left(\alpha\ D_r\ \frac{\partial A_z}{\partial z}\right) - i\ \frac{\gamma_0}{\gamma}\ \sqrt{\alpha}\ D_r k_y\ \frac{\partial A_z}{\partial z} - (\gamma + k_y^2 D_r)A_z = 0. \tag{30}$$

Prior to solving (30) for the density profile shown in Fig. 1b, we first consider a blob of infinite extent along the field. We Fourier expand modes parallel to B_0, i.e., $\tilde{p}(z) \sim \tilde{p}\ \exp(ik_z z)$ and use the local approximation $[k_y(\partial\ \ell n\ n_0/\partial x)^{-1} >> 1]$ so that (30) can be solved

algebraically. This allows comparison with previous results [Ossakow and Chaturvedi, 1979; Chaturvedi and Ossakow, 1981]. The local dispersion equation is given by

$$\gamma - \frac{\gamma_0}{\gamma} k_y k_z \sqrt{\alpha}\, D_r + \left(k_y^2 + \alpha k_z^2\right) D_r = 0. \tag{31}$$

The first term in (31) is related to electromagnetic effects, the second term to convection along the gradient (which causes the instability), and the final term to perpendicular ion motion and parallel electron motion. In the electrostatic limit $\left(k_y^2 D_r,\ \alpha k_z^2 D_r >> \gamma\right)$ the growth rate is given by

$$\gamma = + \gamma_0 \frac{\sqrt{\alpha}\ \left(k_z/k_y\right)}{1 + \alpha k_z^2/k_y^2} \tag{32}$$

which agrees with the results of Ossakow and Chaturvedi (1979). This growth rate has a maximum value when $k_z^2/k_y^2 = 1/\alpha$. It is given by

$$\gamma_{max} = \gamma_0/2 \tag{33}$$

Retention of electromagnetic effects yields the following growth rate

$$\gamma = - \frac{1}{2} \hat{k}^2 D_r \pm \frac{1}{2} \left[\hat{k}^2 D_r + 4\gamma_0 k_y k_z \sqrt{\alpha}\ D_r\right]^{1/2} \tag{34}$$

where $\hat{k}^2 = k_y^2 + \alpha k_z^2$ and which can be shown to agree with Chaturvedi and Ossakow (1981).

We now solve (30) for a density profile relevant to a plasma "blob" in the auroral ionosphere. The equation is solved subject to the boundary

condition $\hat{A}_z \to 0$ as $|z| \to \infty$. For a "top hat" density profile (Fig. 1b), the solution to (30) in the region $z < 0$ and $z > z_0$ can be written as a plane wave,

$$\tilde{A}_z = \tilde{A}_i \exp(-k_i|z|) \tag{35}$$

where

$$k_i = k_y \left(\frac{1}{\sqrt{\alpha R}}\right)_i \tag{36}$$

and

$$R = \frac{k_y^2 D_r}{\gamma + k_y^2 D_r}. \tag{37}$$

Note that the subscript i denotes quantities that are evaluated in the background ionosphere (i.e., $z < 0$ and $z > z_0$). The parameter R is a measure of the electrostatic/electromagnetic nature of the mode. For $\gamma \ll k_y^2 D_r$ we note that $R \sim 1$ and the mode is essentially electrostatic. In the opposite limit, $\gamma \gg k_y^2 D_r$ the mode is essentially electromagnetic with $R \ll 1$.

In the region $0 < z < z_0$, i.e., region occupied by the density enhancement, the solution to $\tilde{A}_z$ is given by

$$\tilde{A}_z = \tilde{A}_{zb}^1 \exp(ik_1 z) + \tilde{A}_{zb}^2 \exp(ik_2 z) \tag{38}$$

where

$$k_1 = \frac{k_y}{\sqrt{\alpha}} \left[-\frac{1}{2} \frac{\gamma_0}{\gamma} + \frac{1}{2} \left(\frac{\gamma_0^2}{\gamma^2} - \frac{4}{R} \right)^{1/2} \right]_b \tag{39}$$

and

$$k_2 = \frac{k_y}{\sqrt{\alpha}} \left[-\frac{1}{2} \frac{\gamma_0}{\gamma} - \frac{1}{2} \left(\frac{\gamma_0^2}{\gamma^2} - \frac{4}{R} \right)^{1/2} \right]_b \tag{40}$$

Again, the subscript b denotes quantities that are to be evaluated in the blob ($0 < z < z_0$).

To determine the dispersion equation we match the plane wave solutions at $z = 0$ and $z = z_0$. The appropriate matching conditions are obtained from (28) and (29). Namely, we require that $\hat{A}_z$ and $\alpha D_r \partial A_z / \partial z$ be continous across each boundary. Using these boundary conditions, we find that the dispersion equation is given by

$$\exp(i \Delta k z_0) = \Lambda \tag{41}$$

where

$$\Lambda = \frac{1 + i\Gamma \frac{k_2}{k_i}}{1 + i\Gamma \frac{k_1}{k_i}} \; \frac{1 - i\Gamma \frac{k_1}{k_i}}{1 - i\Gamma \frac{k_2}{k_i}} \tag{42}$$

$$\Delta k = k_1 - k_2 = \frac{k_y}{\sqrt{\alpha}} \left[\frac{\gamma_0^2}{\gamma^2} - \frac{4}{R} \right]_b^{1/2} \tag{43}$$

$$\Gamma = \frac{(\alpha D_r)_b}{(\alpha D_r)_i} \tag{44}$$

and k_i, k_1, and k_2 are given in (36), (39), and (40), respectively.

IV. ANALYTICAL AND NUMERICAL RESULTS

We now present analytical and numerical results for the growth rate of the current convective instability based upon (41). We cast (41) into dimensionless form and obtain

$$\exp(i\hat{k}_y\hat{z}_0\Delta\hat{k}) = \Lambda \tag{45}$$

where $\hat{k}_y = k_yL_r$, $\hat{z}_0 = z_0/\sqrt{\alpha}\,L_r$, $\Delta\hat{k} = \Delta k/k_y$, and $L_r = D_r/\gamma_0$ is the resistive diffusion length. First, we note that in the short wavelength limit, i.e., when $\hat{k}_y\hat{z}_0 >> 1$, $k_1 = k_2$ which implies $\Delta k = 0$, so that $\Lambda = 1$ is a solution to (45). The growth rate is given by

$$\gamma = \frac{\gamma_0}{2}\sqrt{R_b}. \tag{46}$$

In the electrostatic limit $(\gamma << k_y^2D_r)$ we note that $R \sim 1$ and $\gamma = \gamma_0/2$ which corresponds to the maximum growth rate of the instability based upon local theory [see (33)]. This result makes sense physically since the limit $\hat{k}_y\hat{z}_0 >> 1$ corresponds to modes with parallel wavelengths much smaller than the parallel length of the blob; hence, the outer regions (i.e., $z < 0$ and $z > z_0$) have little affect on the dynamics of the instability.

On the other hand, for long wavelengths, i.e., $\hat{k}_y\hat{z}_0 << 1$, the growth rate is strongly affected by the finite length of the blob. The growth rate is reduced and $k_2 >> k_1$. In this limit $\Lambda = -1$ and the dispersion equation becomes

$$\hat{k}_y\hat{z}_0\,\Delta\hat{k} = (2m + 1)\pi \tag{47}$$

where m = 1, 2, ... is an integer. Equation (47) yields the following growth rate

$$\gamma = \gamma_0 \frac{\hat{k}_y \hat{z}_0}{(2m + 1)\pi} \tag{48}$$

We note that this growth rate is independent of the value of R. Also, the fastest growing mode in this regime corresponds to m = 0.

In Fig. 2 we present numerical results based upon (45). We plot $\hat{\gamma} = \gamma/\gamma_0$ versus $\hat{k}_y = k_y L_r$ for $\hat{z}_0 = z_0/\sqrt{\alpha}\, L_r = 1$ and $\Gamma = 0.2$ for several values of mode number m. The solid curves are numerical results while the dashed curves are based upon the analytical results (46) and (48). Note the excellent agreement between these results in the short wavelength (i.e., $\hat{k}_y >> 2\pi$) and long wavelength (i.e., $\hat{k}_y << 2\pi$) limits. In the short wavelength regime the modes are electrostatic ($R \sim 1$) and have a growth rate which is the same as the maximum value obtained from local theory [$\gamma \sim \gamma_0/2$, see (33)]. In the long wavelength limit the modes are electromagnetic ($R << 1$) and have a smaller growth rate than in the short wavelength limit.

Physically, the reduction in growth can be explained as follows. The current convective instability requires a density and potential perturbation parallel to $\underline{B}_0$ which corresponds to a parallel wavenumber in the local analysis (the second term in (31) is the driver and depends upon γ_0 and k_z). From local theory, the instability achieves maximum growth when $k_z/k_y = 1/\sqrt{\alpha} = (\sigma_\perp/\sigma_\parallel)^{1/2}$ where $\sigma_\perp$ and $\sigma_\parallel$ are the perpendicular and parallel conductivities, respectively. In the ionosphere, $\sigma_\perp/\sigma_\parallel \sim 10^{-8} <<$ 1 so that the instability favors very long wavelengths parallel

to $\underline{B}_0$ relative to the perpendicular wavelengths $(\lambda_\parallel >> \lambda_\perp)$. We define the maximum effective parallel wavenumber that can be supported within the blob as

$$k_{z(eff)} = 2\pi/z_0. \tag{49}$$

Using (49) to determine the value of $\hat{k}_y\hat{z}_0$ for maximum growth based on local theory $(k_y/k_z = \sqrt{\alpha})$ we find that $\hat{k}_y\hat{z}_0 = 2\pi$. Thus, for $\hat{k}_y\hat{z}_0 >> 2\pi$ we expect that the blob is sufficiently long to support a parallel wavelength which yields maximum growth, i.e., $k_z > k_{z(eff)}$. This is, in fact, quite apparent in Fig. 2 where $\hat{\gamma} \sim 0.5$ for $\hat{k}_y\hat{z}_0 >> 2\pi$. However, for modes such that $\hat{k}_y\hat{z}_0 << 2\pi$ it is not possible for any mode to satisfy the condition $k_z/k_y \sim 1/\sqrt{\alpha}$. Therefore, the modes grow at a reduced growth rate; again this is apparent from Fig. 2. In fact, this corresponds to the limit $\sqrt{\alpha}\ k_z/k_y >> 1$ based on local theory. Making use of this approximation in (32) we find that

$$\gamma = \gamma_0 \frac{k_y}{\sqrt{\alpha}\ k_z} \tag{50}$$

which agrees with (48) if we make the identification $k_z^{-1} = z_0/(2m + 1)\pi$.

V. DISCUSSION

We have studied the effects of finite size of high latitude "blobs" parallel to magnetic field on the current convective instability in auroral ionosphere. We find that the growth rate of the instability can be substantially reduced in the case of a finite size "blob" from the value obtained when the "blobs" are assumed to be infinite. For short

wavelengths, such that $k_y z_0 >> \sqrt{\alpha}$ (where k is the wavenumber, z_0 the blob size and $\alpha = \Omega_e \Omega_i / \nu_e \nu_i = \sigma_\parallel / \sigma_\perp$), the growth rate remains relatively unaffected (especially in the electrostatic case) by the blob size. In the long wavelength case, when $k_y z_0 << \sqrt{\alpha}$, the growth rate becomes proportional to the wavenumber and the blob size. Thus, for a given blob size, the longer transverse wavelengths have smaller growth rates and, similarly, for shorter blob sizes, the growth rates are reduced [see (48)].

The current convective instability has been discussed in the auroral ionosphere recently, where field-aligned currents are a constant feature and plasma density enhancements ("blobs") have been observed with structured walls. We now apply our results for the instability to this situation. The typical parameters in this situation are the following: ambient density $n_1 \sim 10^6$ cm^{-3}, blob density $n_0 \sim 5 \times 10^6$ cm^{-3}, field-aligned current density $J_{\parallel 0} \lesssim 10^2$ μamp/m^2 (which corresponds to an electron parallel drift $v_{0\parallel} \lesssim 500$ m/sec for $n_0 \sim 10^6$ cm^{-3}), the scale length associated with the "blob" gradient $L_N \sim 50$ kms, the ion-neutral collision frequency $\nu_{in} \sim 5 \times 10^{-2}$ sec^{-1} and electron collision frequency $\nu_e \sim 5 \times 10^2$ sec $^{-1}$ [Tsunoda and Vickrey, 1985]. For $B_0 \sim 0.5$ G, one has $\nu_e/\Omega_e \sim 6 \times 10^{-5}$ and $\nu_{in}/\Omega_i \sim 1.7 \times 10^{-4}$ so that $\alpha \sim (\Omega_e \Omega_i / \nu_e \nu_{in}) \sim 10^8$. The maximum local growth rate is $\gamma_m \sim \gamma_0/2 \sim 8.5 \times 10^{-3}$ sec^{-1}, where $\gamma_0 \sim (v_{0\parallel}/L_N)(\Omega_e \nu_i / \nu_e \Omega_i) \sim 1.7 \times 10^{-2}$ sec^{-1}. The "blob" dimension parallel to the field is $z_0 \sim 300$ km [Tsunoda and Vickrey, 1985] and for the above set of parameters we have $\omega_{pe} \sim 5.6 \times 10^7$ sec^{-1}, $D_{ri} \sim 1.4 \times 10^8$ cm^2/sec, $L_r \sim D_{ri}/\gamma_0 \sim 8 \times 10^9$ cm, and $\hat{z}_0 \sim z_0/\sqrt{\alpha}\, L_r \sim 3.8 \times 10^{-7}$. For $\lambda_\perp \sim 1$ km, $\hat{k}_y \sim k_y L_r \sim 5 \times 10^5$, and $\hat{k}_y \hat{z}_0 \sim 0.2$. We find from (48), that the nonlocal growth rate in this instance is $\gamma = 0.06\ \gamma_0 = 1.0 \times 10^{-3}$ sec^{-1} for the fastest growing mode (m = 0). This is an order of magnitude smaller

than the maximum local growth γ_m, i.e., $\gamma_m = 8.5 \times 10^{-3}$ sec^{-1}. For a similar set of parameters, we find that for $\lambda_\perp \sim 10$ km, the reduction in the growth rate is even more substantial, i.e., $\gamma = 0.006\ \gamma_0 = 1.0 \times 10^{-4}$ sec^{-1}. Therefore, the current convective instability growth rates may be too small for the instability to explain the blob-associated structure for irregularity scale-sizes of 1 - 10 km.

It may be pointed out that the finite blob-size induced reduction in the current convective instability growth rate of the instability is also related to our assumption that the instability region is confined to the "blob" (i.e., the driving density gradient is negligible outside the blob, namely, in the ambient ionosphere). This is represented by the requirement that the modes decay exponentially outside the blob. This is a reasonable assumption in the blob-associated structure studies. However, should there be a transverse gradient in the ambient ionosphere at high altitude F-region, and if a positive correlation is observed between the structure and the field-aligned currents, then the current convective instability could be considered a possible mechanism for scale sizes on the order of a few km or less.

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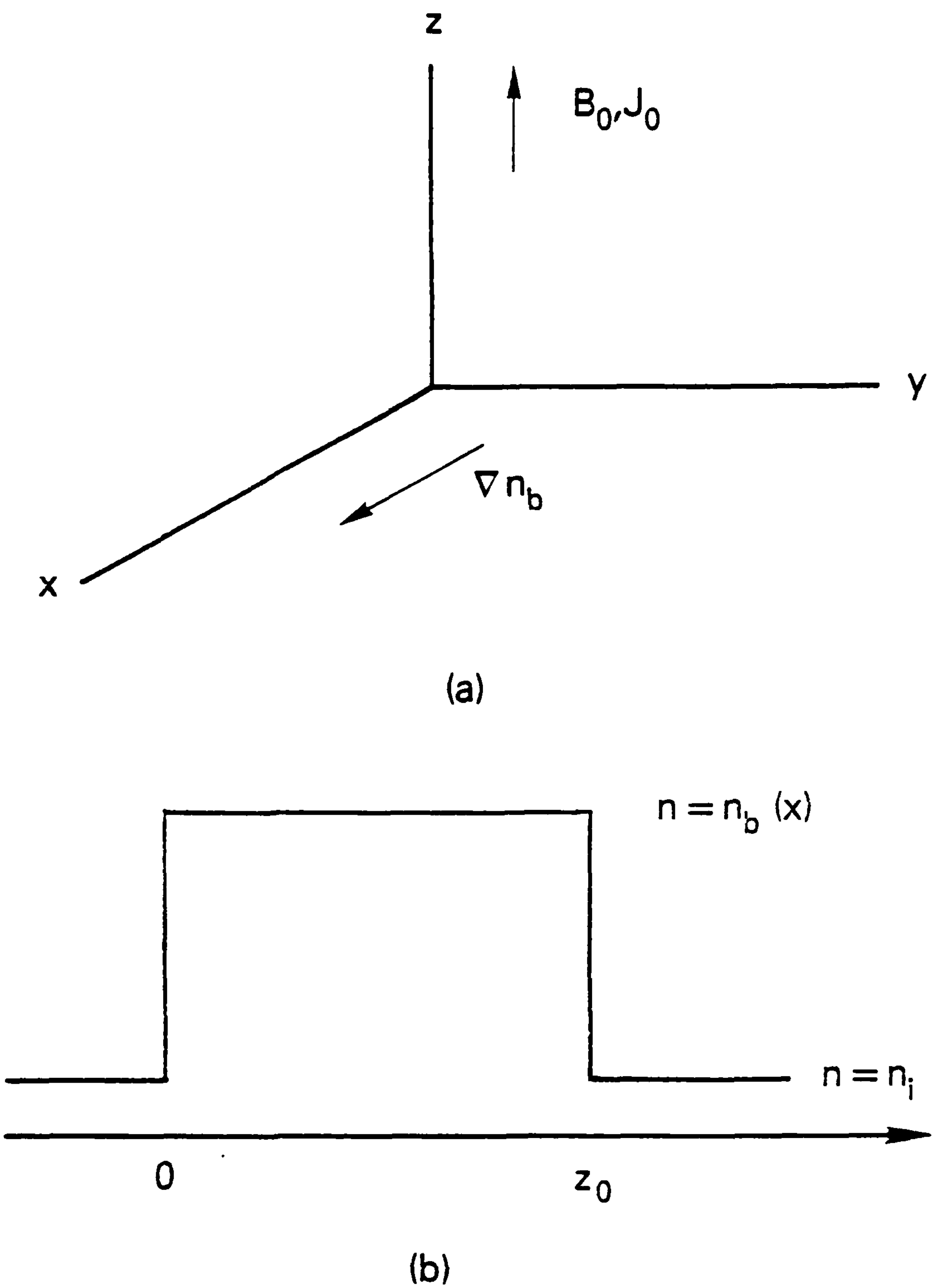

Fig. 1) Plasma configuration and slab geometry. (a) Coordinate system and ambient plasma properties. (b) "Top hat" density model for plasma along the ambient magnetic field.

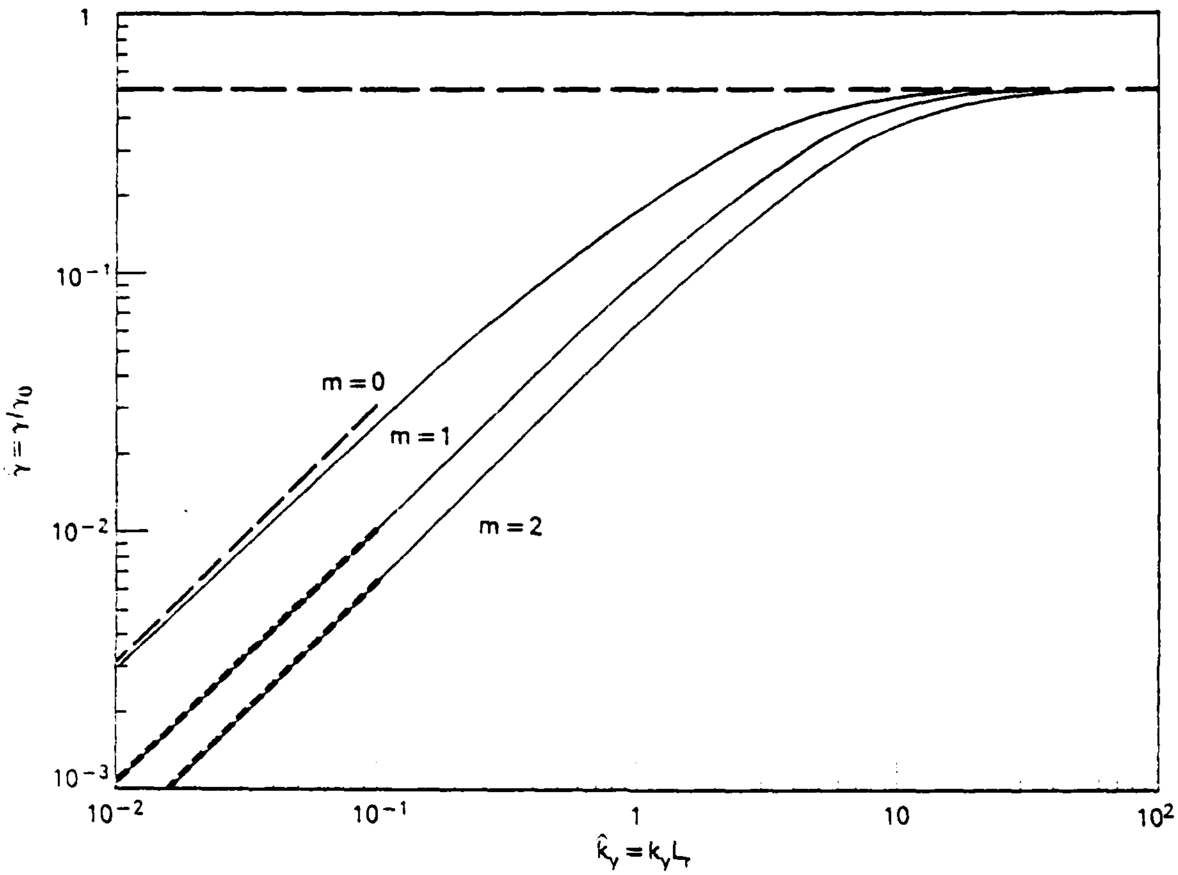

Fig. 2) Plot of growth rate $(\hat{\gamma} = \gamma/\gamma_0)$ versus perpendicular wavenumber $(\hat{k}_y = k_y L_r)$ for several mode numbers ($m = 0,1,2$). The solid curves are numerical solutions of the dispersion equation (45) while the dashed curves are analytic solutions. For $\hat{k}_y \gg 1$ we use (46) while for $\hat{k}_y \ll 1$ we use (48).

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